

Nonminimal couplings to gravity and axion isocurvature bounds

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For axions present during inflation, it has been shown that a nonminimal coupling ξ_σ of the inflaton to gravity worsens isocurvature bounds [arXiv:2504.02952], while a nonminimal coupling ξ_ρ of the radial Peccei-Quinn field can alleviate them [Revisiting isocurvature bounds on the minimal QCD axion, *J. High Energy Phys.* **12** (2025) 028.]. We analyze the simultaneous presence of both couplings and determine when one effect dominates the other, in both the metric and Palatini formulations of gravity. The two tendencies interpolate smoothly, but introducing a nonminimal inflaton coupling reduces the viable interval of ξ_ρ in which isocurvature bounds can be alleviated while avoiding backreaction on the inflationary dynamics. We illustrate our findings in Palatini Higgs inflation and Starobinsky inflation.

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I. INTRODUCTION

Axions [1–3] are among the best-motivated proposals for physics beyond the Standard Model. They provide a viable dark matter candidate and, moreover, offer a solution to the strong CP problem¹ (see, e.g., the review [9]). In the original proposal [1], axions arise as angular Goldstone bosons associated with the breaking of a global Peccei-Quinn (PQ) symmetry. While other ultraviolet completions have been proposed—including constructions based on extra dimensions [10] (see also [11]), local gauge invariance [8,12–14], and Einstein-Cartan gravity [15–17]—in this work we focus on PQ axions.

The cosmological evolution of axions depends on whether the PQ symmetry is broken before or after inflation. If the symmetry is broken after inflation, topological defects form and may overclose the universe in certain scenarios [18]. If PQ symmetry is broken before inflation (and is not restored afterward), the presence of a massless axion field during inflation induces

isocurvature fluctuations [19]. Their nonobservation, especially in cosmic microwave background (CMB) measurements [20], imposes stringent constraints on the compatibility of preinflationary axions with inflation. Avoiding these isocurvature bounds generally requires a large axion decay constant f_a or a low inflationary Hubble scale H .

Some of the leading inflationary models favored by current CMB data [20,21]—such as Starobinsky inflation [22], the metric [23] and Palatini [24] versions of Higgs inflation, and some classes of attractor models [25,26]²—share a common structural feature: they all involve a nonminimal coupling of the inflaton σ to gravity of the form $\xi_\sigma \sigma^2 R$, where R is the Ricci scalar and $\xi_\sigma \gg 1$ a corresponding coupling constant. Recently, however, we pointed out that such a nonminimal coupling inevitably reduces the inflationary value $f_a^{(\text{inf})}$ of the decay constant [29],

$$f_a^{(\text{inf})} = \frac{f_a}{\sqrt{1 + \xi_\sigma \frac{\sigma^2}{M_P^2}}}, \quad (1)$$

where M_P is the Planck mass.

²The main feature of α -attractor models, a pole in the noncanonical kinetic term, can be obtained from a negative nonminimal coupling, $\xi_\rho < 0$ [27,28]. In this case, it can be possible to alleviate isocurvature bounds [29]. However, α -attractor models can also be derived independently of a nonminimal coupling [27,30].

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¹There are indications that quantum gravity does not tolerate eternal de Sitter states due to a fundamental inconsistency caused by quantum breaking [4–6]. This turns the existence of an axion from a naturalness question into a consistency requirement [7,8].

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It has long been known [31–36] that a reduced value $f_a^{(\text{inf})} < f_a$ strengthens isocurvature constraints, while only $f_a^{(\text{inf})} > f_a$ can relax them.³ Contrary to earlier claims [47], this implies that Palatini Higgs inflation is incompatible with isocurvature bounds [29], while for Starobinsky inflation and metric Higgs inflation a comparable tension already exists even without considering the effect of the modified decay constant.

On the other hand, one may also introduce a nonminimal coupling $\xi_\rho \rho^2 R$ for the radial component ρ of the PQ field [48–50], with coupling constant ξ_ρ . It was recently shown that such a coupling can increase $f_a^{(\text{inf})}$ and thereby relax isocurvature bounds [51].⁴ In the relevant parameter regime, the inflationary decay constant takes the approximate form [51]

$$f_a^{(\text{inf})} \approx \sqrt{\frac{12\xi_\rho}{\lambda_\rho}} H, \quad (2)$$

where λ_ρ is the self-coupling of ρ . A sufficiently large ξ_ρ and sufficiently small λ_ρ can therefore yield $f_a^{(\text{inf})} > f_a$, alleviating isocurvature constraints.

In this paper, we *simultaneously* consider the inflaton coupling ξ_σ and the PQ-field coupling ξ_ρ , thus combining the strengthening of isocurvature bounds induced by ξ_σ [29] with the relaxation enabled by ξ_ρ [51]. In brief, we find that these effects interpolate smoothly: For sufficiently small ξ_ρ , the nonminimal inflaton coupling worsens isocurvature bounds as in (1), whereas for sufficiently large ξ_ρ , the inflationary decay constant becomes insensitive to ξ_σ and the result (2) is recovered. However, the parameter space in which (2) applies depends on ξ_σ . While the lower bound on ξ_ρ above which isocurvature constraints can be alleviated is essentially independent of ξ_σ , a nonminimal inflaton coupling reduces the upper bound on ξ_ρ . This limits the maximal enhancement of the inflationary decay constant. As we shall demonstrate, the restriction on ξ_ρ arises primarily from the requirement that the nonminimally coupled PQ field must not interfere with the dynamics of inflation.

Whereas the result of [51] was obtained in the metric formulation of general relativity (GR), we extend the analysis to Palatini gravity, showing that the choice of formulation does not significantly affect the outcome. For each formulation, we select a representative inflationary model for detailed study: Palatini Higgs inflation [23,24], following [29], and the Starobinsky model [22]. Besides

³See [32,37–46] for further approaches to evading isocurvature bounds.

⁴Furthermore, a nonminimal coupling ξ_ρ can be used to drive inflation with the radial PQ-field [35,36,52–54] (see also [45,55,56]).

warm inflation (see [57] and [58–62]), [51] also considered Starobinsky inflation, though without accounting for the effect of ξ_σ on $f_a^{(\text{inf})}$. We find that the mechanism proposed in [51] for alleviating isocurvature constraints indeed applies to Starobinsky inflation, but only within a smaller region of parameter space than was considered in [51].

The paper is organized as follows. In Sec. II, we introduce the model featuring both the axion and the inflaton coupled nonminimally to gravity. We show that the decay constant generically depends on the nonminimal couplings and discuss the resulting impact on isocurvature constraints. In Sec. III, we analyze the limiting case where only the axion couples nonminimally to gravity, i.e., $\xi_\sigma = 0$. Finally, Sec. IV presents the complete analysis including both nonminimal couplings, in both the Palatini and metric formulations of GR. As a key example, we demonstrate that it is possible to satisfy isocurvature bounds in Starobinsky inflation, albeit within a restricted region of parameter space.

Convention: We use the metric signature $(-1, 1, 1, 1)$.

II. THE MODEL

We consider an inflaton field σ with nonminimal coupling ξ_σ and the PQ-field Φ with nonminimal coupling ξ_ρ ⁵

$$\mathcal{L} = \left(\frac{M_P^2}{2} + \xi_\rho |\Phi|^2 + \frac{1}{2} \xi_\sigma \sigma^2 \right) R - \partial_\mu \Phi \partial^\mu \Phi^\star - V_\rho - \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - V_\sigma, \quad (3)$$

where we assume $\xi_\sigma > 0$, $\xi_\rho > 0$ and

$$V_\rho = \lambda_\rho \left(|\Phi|^2 - \frac{1}{2} f_a^2 \right)^2. \quad (4)$$

Splitting the PQ field into

$$\Phi = \frac{\rho}{\sqrt{2}} \exp(i\varphi), \quad (5)$$

one recognizes the nonminimal coupling to the scalar curvature

$$\Omega^2 = 1 + \frac{\xi_\rho \rho^2}{M_P^2} + \frac{\xi_\sigma \sigma^2}{M_P^2}. \quad (6)$$

We perform the conformal transformation $g_{\mu\nu} \rightarrow \Omega^{-2} g_{\mu\nu}$, to obtain (see [63])

⁵Note that these terms are allowed by symmetry, and of mass dimension less or equal to 4, so one should generically include them.

$$\begin{aligned}\mathcal{L} = & \frac{M_P^2}{2}R - \frac{\rho^2}{2\Omega^2}\partial_\mu\varphi\partial^\mu\varphi - \frac{1}{2\Omega^2}\left(1 + \frac{\zeta 6\xi_\rho^2\rho^2}{M_P^2\Omega^2}\right)\partial_\mu\rho\partial^\mu\rho \\ & - \frac{V_\rho}{\Omega^4} - \frac{1}{2\Omega^2}\left(1 + \frac{\zeta 6\xi_\sigma^2\sigma^2}{M_P^2\Omega^2}\right)\partial_\mu\sigma\partial^\mu\sigma - \frac{V_\sigma}{\Omega^4} \\ & - \frac{1}{\Omega^2}\frac{\zeta 6\xi_\rho\xi_\sigma\rho\sigma}{M_P^2\Omega^2}\partial_\mu\rho\partial^\mu\sigma.\end{aligned}\quad (7)$$

As is well known, the outcome of the conformal transformation depends on the formulation of GR. In order to account for this fact, we introduced the parameter ζ , where $\zeta = 1$ for metric GR while $\zeta = 0$ in the Palatini case (see, e.g., [63] for more details on the equivalent formulations of GR).

A. Inflationary decay constant

From the first line of Eq. (7), we see that the inflationary value of the decay constant is given by

$$f_a^{(\text{inf})} = \frac{\rho_{\min}}{\Omega}, \quad (8)$$

where $\rho_{\min}$ is the field value that minimizes the potential of field ρ for a given inflationary background. It is important to note that $f_a^{(\text{inf})}$ is independent of the noncanonical kinetic term of the radial field ρ .⁶ As shown in [29], we immediately see why a nonminimal coupling of the inflaton worsens isocurvature bounds: If $\xi_\rho = 0$, we get $\rho_{\min} = f_a$ and so $\Omega \gg 1$ decreases the inflationary decay constant [29], as shown in Eq. (1). What remains to be done is to evaluate $\rho_{\min}$ for nonvanishing ξ_ρ , which we shall do in different scenarios.

B. Isocurvature bounds

Before that, we will briefly state the known isocurvature bound [19]. Assuming that axions make up all of dark matter, it reads [29,33–35,51,54]

$$\frac{H}{2\pi f_a^{(\text{inf})}} \lesssim 4.6 \times 10^{-6} \left(\frac{1.02 \times 10^{12} \text{ GeV}}{f_a} \right)^{7/12}. \quad (9)$$

Plugging in the result (2) of [51] for $f_a^{(\text{inf})}$ then gives

$$\lambda_\rho < 10^{-8} \xi_\rho \left(\frac{1.02 \times 10^{12} \text{ GeV}}{f_a} \right)^{7/6}. \quad (10)$$

Thus, evading isocurvature bounds generically requires a tiny λ_ρ or a large $\xi_\rho \gg 1$.

⁶This is because we are interested in the minimum of the potential, which happens when the kinetic term vanishes $\partial_\mu\rho = 0$.

III. INFLATON MINIMAL COUPLING

First, we shall consider the case of a minimally coupled inflaton, $\xi_\sigma = 0$, as in [51]. The following derivation applies both to the metric and Palatini formulations. Since there is no kinetic mixing between ρ and σ , Eq. (7) directly gives the effective potential of the two scalar fields

$$U = \frac{V_\rho + V_\sigma}{\Omega^4} = \frac{\frac{\lambda_\rho}{4}(\rho^2 - f_a^2)^2 + V_\sigma}{\left(1 + \frac{\xi_\rho\rho^2}{M_P^2}\right)^2}. \quad (11)$$

Minimizing U , we get

$$\rho_{\min} = f_a \frac{\sqrt{1 + \frac{\xi_\rho f_a^2}{M_P^2} + \frac{4\xi_\rho V_\sigma}{\lambda_\rho f_a^2 M_P^2}}}{\sqrt{1 + \frac{\xi_\rho f_a^2}{M_P^2}}}, \quad (12)$$

and plugging this into Eq. (6) yields

$$\Omega^2 = 1 + \frac{\xi_\rho f_a^2}{M_P^2} + \frac{4\xi_\rho^2 V_\sigma}{\lambda_\rho M_P^4 \left(1 + \frac{\xi_\rho f_a^2}{M_P^2}\right)}. \quad (13)$$

This determines the inflationary decay constant via Eq. (8).

As evident from the third line of Eq. (7), the presence of a nonminimal coupling ξ_ρ of the PQ-field leads to a non-trivial coefficient $1/\Omega^2$ of the inflaton kinetic term. In order to avoid changing inflation, we need $\Omega^2 \approx 1$. So $\Omega^2 - 1 \ll 1$ leads to the necessary requirement

$$\xi_\rho \ll \min\left(\frac{M_P^2}{f_a^2}, \sqrt{\lambda_\rho} \frac{M_P^2}{\sqrt{V_\sigma}}\right), \quad (14)$$

where the second part coincides with its counterpart derived in [51]. In general, however, the condition $U \approx V_\sigma$ may not be sufficient for leaving the dynamics of the inflaton unaltered. In order to check this, we can evaluate the correction to the first slow-roll parameter

$$\frac{\sqrt{\epsilon} - \sqrt{\epsilon}|_{\xi_\rho=0}}{\sqrt{\epsilon}|_{\xi_\rho=0}} \approx \Omega - 1 \sim \frac{\xi_\rho \rho_{\min}^2}{M_P^2}, \quad (15)$$

where we used that the canonical inflaton field χ satisfies approximately $d\chi \approx d\sigma/\Omega$. We reproduce the condition $\Omega - 1 \ll 1$, and so for the case of a minimally coupled inflaton, the condition (14) is both sufficient and necessary for not changing ϵ .

If ξ_ρ obeys the bound (14), plugging Eq. (12) into the potential (11) shows that $U(\rho_{\min}) \approx V_\sigma$. So $V_\sigma \approx 3M_P^2 H^2$ and we can approximate

$$f_a^{(\text{inf})} \approx \rho_{\min} \approx f_a \left(\sqrt{1 + \frac{12\xi_\rho H^2}{\lambda_\rho f_a^2}} \right). \quad (16)$$

As it should, the value (16) of the inflationary decay constant coincides with the result (2) of [51], which has been derived in the Jordan frame by minimizing $V_{\text{eff}} \approx V_\rho - \xi_\rho R |\Phi|^2 \approx V_\rho - 12\xi_\rho H^2 |\Phi|^2$. We have rederived this finding in the Einstein frame and shown that it is also applicable in the Palatini formulation of GR.

In summary, the mechanism for lifting isocurvature constraints is effective if⁷

$$\begin{aligned} \frac{\lambda_\rho f_a^2}{H^2} &\ll \xi_\rho \ll \sqrt{\lambda_\rho} \frac{M_P}{H} \\ \Rightarrow f_a^{(\text{inf})} &\approx \sqrt{\frac{12\xi_\rho}{\lambda_\rho}} H \gtrsim f_a, \end{aligned} \quad (17)$$

where we took into account Eqs. (14) and (16). Such a ξ_ρ only exists if

$$\lambda_\rho f_a^4 \ll M_P^2 H^2. \quad (18)$$

In general, the maximal temperature during reheating, i.e., the transition from inflation to radiation-dominated expansion, fulfills (see, e.g., [64,65])

$$T_{\text{max}} < \left(\frac{90 M_P^2 H^2}{\pi^2 g_\star} \right)^{1/4}. \quad (19)$$

This bound is saturated if reheating can be approximated as instantaneous. In this case, the hierarchy (18) and the condition $T_{\text{max}} < f_a$, which is necessary to avoid restoration of the PQ-symmetry, can only be satisfied simultaneously if $\lambda_\rho \ll (T_{\text{max}}/f_a)^4 \ll 1$.

IV. GENERAL CASE: FULL NONMINIMAL COUPLINGS

A. Palatini gravity

Next, we shall include a nonminimal coupling ξ_σ of the inflaton, but specialize to the Palatini formulation of GR. Correspondingly, we set $\zeta = 0$ in the Einstein frame action (7). Since no kinetic mixing of the two scalar fields exists in Palatini GR, the effective potential becomes

$$U = \frac{V_\sigma + V_\rho}{\Omega^4} = \frac{\frac{\lambda_\rho}{4}(\rho^2 - f_a^2)^2 + V_\sigma}{\left(1 + \frac{\xi_\rho \rho^2}{M_P^2} + \frac{\xi_\sigma \sigma^2}{M_P^2}\right)^2}, \quad (20)$$

⁷Note that $\lambda_\rho f_a^2/H^2 \ll \xi_\rho \ll \sqrt{\lambda_\rho} M_P/H$ implies $\sqrt{\lambda_\rho} M_P/H \gg M_P^2/f_a^2$, and so we only need to consider the second condition of Eq. (14).

which is minimized for

$$\rho_{\min} = f_a \sqrt{\frac{1 + \frac{\xi_\rho f_a^2}{M_P^2} + \frac{\xi_\sigma \sigma^2}{M_P^2} + \frac{4\xi_\rho V_\sigma}{\lambda_\rho f_a^2 M_P^2}}{1 + \frac{\xi_\rho f_a^2}{M_P^2} + \frac{\xi_\sigma \sigma^2}{M_P^2}}}, \quad (21)$$

and so

$$\Omega^2 = \frac{\left(1 + \frac{\xi_\rho f_a^2}{M_P^2} + \frac{\xi_\sigma \sigma^2}{M_P^2}\right)^2 + \frac{4\xi_\rho^2 V_\sigma}{\lambda_\rho M_P^4}}{1 + \frac{\xi_\rho f_a^2}{M_P^2} + \frac{\xi_\sigma \sigma^2}{M_P^2}}. \quad (22)$$

As before, this determines the inflationary decay constant via Eq. (8).

We can now split the full nonminimal coupling Ω into the axion and inflaton dependent parts. Introducing the notation

$$\Omega_0^2 \equiv 1 + \frac{\xi_\sigma \sigma^2}{M_P^2}, \quad (23)$$

we shall assume that (already in the absence of an axion) $\Omega_0^2 \gg 1$. Thus, a necessary condition for avoiding back-reaction on the inflaton is to require $\Omega^2 \approx \Omega_0^2$, and so $\Omega^2 - \Omega_0^2 \ll \Omega_0^2$ implies

$$\xi_\rho \ll \min \left(\Omega_0^2 \frac{M_P^2}{f_a^2}, \sqrt{\lambda_\rho} \frac{M_P}{H} \right), \quad (24)$$

where we self-consistently used that $3M_P^2 H^2 \approx V_\sigma/\Omega_0^4 \approx M_P^4 V_\sigma/(\xi_\sigma^2 \sigma^4)$. The first condition of this bound is less stringent than its counterpart (14) in the minimally coupled case. Provided Eq. (24) is satisfied, we can approximate

$$f_a^{(\text{inf})} \approx f_a \sqrt{\frac{1}{\Omega_0^2} + \frac{12\xi_\rho H^2}{\lambda_\rho f_a^2}}, \quad (25)$$

with the limits

$$f_a^{(\text{inf})} \approx \begin{cases} \frac{f_a}{\Omega_0} & \text{for } \xi_\rho \ll \frac{\lambda_\rho f_a^2}{H^2 \Omega_0^2} \\ \sqrt{\frac{12\xi_\rho}{\lambda_\rho}} H & \text{for } \xi_\rho \gg \frac{\lambda_\rho f_a^2}{H^2 \Omega_0^2}. \end{cases} \quad (26)$$

The first line coincides with the result (8) and so leads to $f_a^{(\text{inf})} \ll f_a$, which results in a worsening of isocurvature bounds. Interestingly, the second line is identical to the finding obtained for a minimally coupled inflaton [see Eqs. (2) and (17)]. Thus, if the coupling to the axion is sufficiently large, a nonminimal coupling of the inflaton no

longer affects the inflationary decay constant. However, comparison of Eqs. (16) and (25) shows that for a given choice of parameters, $f_a^{(\text{inf})}$ is always smaller for a non-minimally coupled inflaton as compared to the minimally coupled case.

In summary, the relevant parameter space for lifting isocurvature constraints is⁸

$$\begin{aligned} \frac{\lambda_\rho f_a^2}{H^2} &\ll \xi_\rho \ll \sqrt{\lambda_\rho} \frac{M_P}{H} \\ \Rightarrow f_a^{(\text{inf})} &\approx \sqrt{\frac{12\xi_\rho}{\lambda_\rho}} H \gtrsim f_a. \end{aligned} \quad (27)$$

It is important to note that $f_a^{(\text{inf})} \approx \sqrt{12\xi_\rho/\lambda_\rho} H$ already holds for smaller $\xi_\rho \gg \lambda_\rho f_a^2/(\Omega_0^2 H^2)$, but achieving $f_a^{(\text{inf})} > f_a$ requires the more stringent lower bound $\xi_\rho \gg \lambda_\rho f_a^2/H^2$. Thus, Eq. (27) coincides with its counterpart (17) of a minimally coupled inflaton. However, the bounds on ξ_ρ of Eq. (17) are only necessary conditions for the validity of lifting isocurvature constraints, but in general they are not sufficient. In particular, the upper bound on ξ_ρ can be significantly stronger than shown in Eq. (27), as we shall demonstrate shortly [see Eq. (29)].

Palatini Higgs inflation. We can come back to isocurvature bounds in Palatini Higgs inflation [23,24], as studied in [29]. Then σ is the Higgs field (in unitary gauge), and so $V = \lambda/4\sigma^4$, where λ is the high-energy value of the Higgs self-coupling. In this concrete model, we can evaluate the first slow-roll parameter and its leading correction

$$\frac{\sqrt{\epsilon} - \sqrt{\epsilon}|_{\xi_\rho=0}}{\sqrt{\epsilon}|_{\xi_\rho=0}} \sim \Omega - 1 \sim \xi_\rho \frac{\rho_{\min}^2}{M_P^2}. \quad (28)$$

Requiring it to be small implies [cf., Eq. (25)]

$$\xi_\rho \ll \min \left(\frac{M_P^2}{f_a^2}, \frac{\sqrt{\lambda_\rho} M_P}{\Omega_0 H} \right). \quad (29)$$

This condition, derived from the first derivative of the potential, is significantly more restricting than the condition (24) derived from the value of the potential itself. Comparing Eq. (29) with its counterpart (14) in the minimally coupled case, we see that the first conditions coincide while the second one is considerably stronger in the nonminimally coupled case.

Matching the amplitude of CMB perturbations requires $\xi_\sigma \sim 10^7$ and during inflation $\sigma \sim \sqrt{N} M_P$ with CMB

⁸For an analogous reason as in footnote 7, we can drop $\Omega_0^2 M_P^2/f_a^2$ in comparison to $\sqrt{\lambda_\rho} M_P/H$.

generation at $N \approx 51$ (see [65]). Consequently, we have $\Omega_0^2 \approx 5 \times 10^8$ and moreover $H \approx 10^{-2} M_P/\xi_\sigma \approx 10^{-9} M_P$. In general, the second condition in Eq. (29) is more stringent and the largest admissible nonminimal coupling becomes $\xi_{\rho,\text{max}} \sim 10^5 \sqrt{\lambda_\rho}$. Plugging this into the bound (10) gives

$$\lambda_\rho < 10^{-6} \left(\frac{10^{12} \text{ GeV}}{f_a} \right)^{7/3}. \quad (30)$$

On the other hand, reheating in Palatini Higgs inflation can be well approximated as instantaneous [64,66], and so Eq. (19) gives $T_{\text{max}} \approx 4 \times 10^{13} \text{ GeV}$, where we took $g_\star \approx 100$. Since Ω_0 changes rapidly during reheating [64,66], the relevant value the axionic decay constant and the precise condition for nonrestoration of PQ-symmetry remain to be determined. Nevertheless, it is reasonable to expect that $T_{\text{max}} > f_a$ will excite the radial mode of the PQ-field (see also [51]). Thus, plugging $f_a > 4 \times 10^{13} \text{ GeV}$ in the bound (30) shows that isocurvature bound can only be lifted at the price of an extremely small $\lambda_\rho \lesssim 10^{-9}$. For different values of λ_ρ , we show in Fig. 1 the influence of the inflationary decay constant (25) on isocurvature bounds.

B. Metric gravity

For metric GR, the situation is more complicated due to the kinetic mixing in the last line of Eq. (7). In order to obtain two at least approximately independent scalar fields, one needs to perform an appropriate shift of σ ,

$$\sigma \rightarrow \sigma + f(\rho, \sigma), \quad (31)$$

where $f(\rho, \sigma)$ depends on both ρ and σ . This generates additional terms in the potential of ρ , which depend on V_σ . As a result, $\rho_{\min}$ might change, although this has to be analyzed on a case by case basis.

Starobinsky inflation. We shall work out one particularly important model of a nonminimally coupled inflaton in metric GR: Starobinsky inflation [22]. The theory is

$$\mathcal{L} = \left(\frac{M_P^2}{2} + \xi_\rho |\Phi|^2 \right) R + \frac{M_P^2}{12M^2} R^2 - \partial_\mu \Phi \partial^\mu \Phi^\star - V_\rho, \quad (32)$$

which replaces Eq. (3). As usual, one can then introduce an auxiliary scalar field σ to replace the R^2 term

$$\begin{aligned} \mathcal{L} = & \left(\frac{M_P^2}{2} + \frac{\xi_\rho \rho^2}{2} + \frac{M_P^2 \sigma^2}{6M^2} \right) R - \frac{M_P^2}{12M^2} \sigma^4 - \frac{1}{2} \partial_\mu \rho \partial^\mu \rho \\ & - \frac{\rho^2}{2} \partial_\mu \varphi \partial^\mu \varphi - V_\rho, \end{aligned} \quad (33)$$

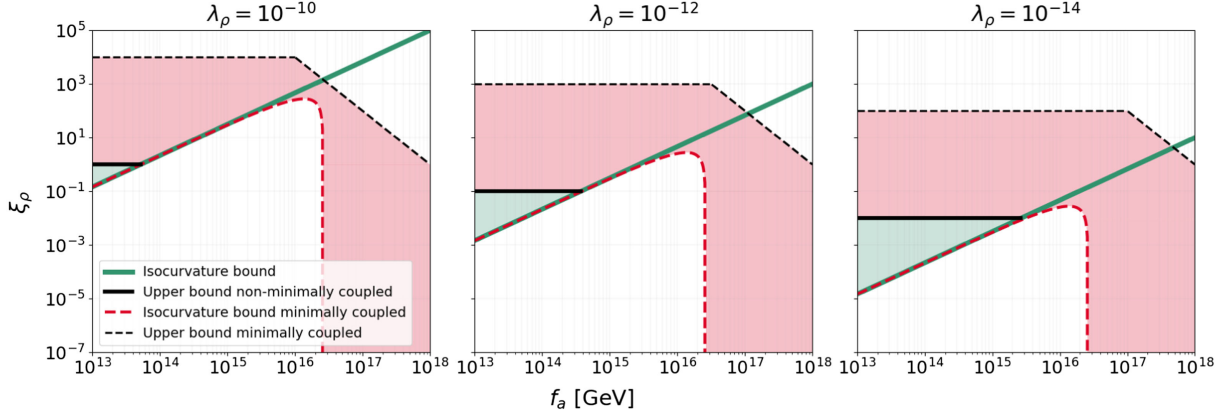


FIG. 1. Constraints on ξ_ρ in Palatini Higgs inflation as a function of f_a , for $H = 10^9$ GeV and some choices of λ_ρ . The green line represents the lower bound on ξ_ρ for which the inflationary decay constant (25) fulfills the isocurvature bound (9). The red curve shows the would-be lower bound on ξ_ρ if the inflaton were minimally coupled, which is found by plugging Eq. (16) into Eq. (9). Furthermore, the solid black line represents the upper bound on ξ_ρ from imposing the nonbackreaction condition of Eq. (29) (with $\Omega_0 \sim 2 \times 10^4$) and the dashed black line corresponds to the would-be upper bound for the minimally coupled case. In the green region, isocurvature bounds are obeyed and the red region would only be viable if the influence of the inflaton nonminimal coupling on axions were neglected. That the red region on the right extends down to vanishing ξ_ρ reflects the fact that Palatini Higgs inflation would obey isocurvature bound if the effect of the inflaton nonminimal coupling were not taken into account [47]. The code used to produce this figure is available at Ref. [67].

where we also plugged in the decomposition (5) of Φ . As before, its presence should not alter inflationary dynamics, which implies

$$\xi_\rho \ll \frac{M_P^2 \sigma^2}{M^2 \rho_{\min}^2}. \quad (34)$$

What is unique about Starobinsky inflation is that the action (33) does not contain a kinetic term for σ , i.e., it is only generated through the conformal transformation. Therefore, it is convenient to perform the shift (31) already in Eq. (33). Redefining

$$\sigma^2 \rightarrow \sigma^2 - \frac{3\xi_\rho \rho^2 M^2}{M_P^2}, \quad (35)$$

we get⁹

$$\begin{aligned} \mathcal{L} = & \left(\frac{M_P^2}{2} + \frac{\sigma^2 M_P^2}{6M^2} \right) R - \frac{M_P^2}{12M^2} \sigma^4 - \frac{1}{2} \partial_\mu \rho \partial^\mu \rho - \frac{\rho^2}{2} \partial_\mu \varphi \partial^\mu \varphi \\ & - V_\rho + \frac{\xi_\rho}{2} \rho^2 \sigma^2 - \frac{3\xi_\rho^2 M^2}{4M_P^2} \rho^4. \end{aligned} \quad (36)$$

The first line of Eq. (36) describes pure Starobinsky inflation in the absence of an axion. Therefore, we

⁹Note that unlike in the previously considered inflationary scenarios, imposing $\Omega^2 - 1 \ll 1$ does not lead to an upper bound on ξ_ρ in terms of M_P^2/f_a^2 [cf., Eqs. (17) and (27)] since the nonminimally coupling $\rho^2 R$ can be removed by the redefinition (35) of the nonpropagating field σ .

can perform the conformal transformation $g_{\mu\nu} \rightarrow \Omega^{-2} g_{\mu\nu}$ with $\Omega^2 = 1 + \sigma^2/(3M^2)$ to obtain

$$\begin{aligned} \mathcal{L} = & \frac{M_P^2}{2} R - \frac{M_P^2}{12M^2} \frac{\sigma^4}{\Omega^4} - \frac{M_P^2 \sigma^2}{3M^4 \Omega^4} \partial_\mu \sigma \partial^\mu \sigma - \frac{1}{2\Omega^2} \partial_\mu \rho \partial^\mu \rho \\ & - \frac{\rho^2}{2\Omega^2} \partial_\mu \varphi \partial^\mu \varphi - \frac{1}{\Omega^4} \left(V_\rho - \frac{\xi_\rho}{2} \rho^2 \sigma^2 + \frac{3\xi_\rho^2 M^2}{4M_P^2} \rho^4 \right). \end{aligned} \quad (37)$$

Introducing the canonically normalized inflaton χ via

$$\sigma^2 = 3M^2 \left(e^{\sqrt{2/3}\chi/M_P} - 1 \right), \quad (38)$$

and using the potential term as defined in Eq. (4), we arrive at

$$\begin{aligned} \mathcal{L} = & \frac{M_P^2}{2} R - \frac{3}{4} M_P^2 M^2 \left(1 - e^{-\sqrt{2/3}\chi/M_P} \right)^2 - \frac{1}{2} \partial_\mu \chi \partial^\mu \chi \\ & - \frac{1}{2e^{\sqrt{2/3}\chi/M_P}} \partial_\mu \rho \partial^\mu \rho - \frac{1}{2e^{\sqrt{2/3}\chi/M_P}} \rho^2 \partial_\mu \varphi \partial^\mu \varphi \\ & - \frac{\lambda_\rho}{4e^{\sqrt{2/3}\chi/M_P}} \left(\left(1 + \frac{3\xi_\rho^2 M^2}{\lambda_\rho M_P^2} \right) \rho^4 \right. \\ & \left. - \left(\frac{6\xi_\rho M^2}{\lambda_\rho} \left(e^{\sqrt{2/3}\chi/M_P} - 1 \right) + 2f_a^2 \right) \rho^2 + f_a^4 \right). \end{aligned} \quad (39)$$

Therefore, we can approximate the potential of ρ during inflation as

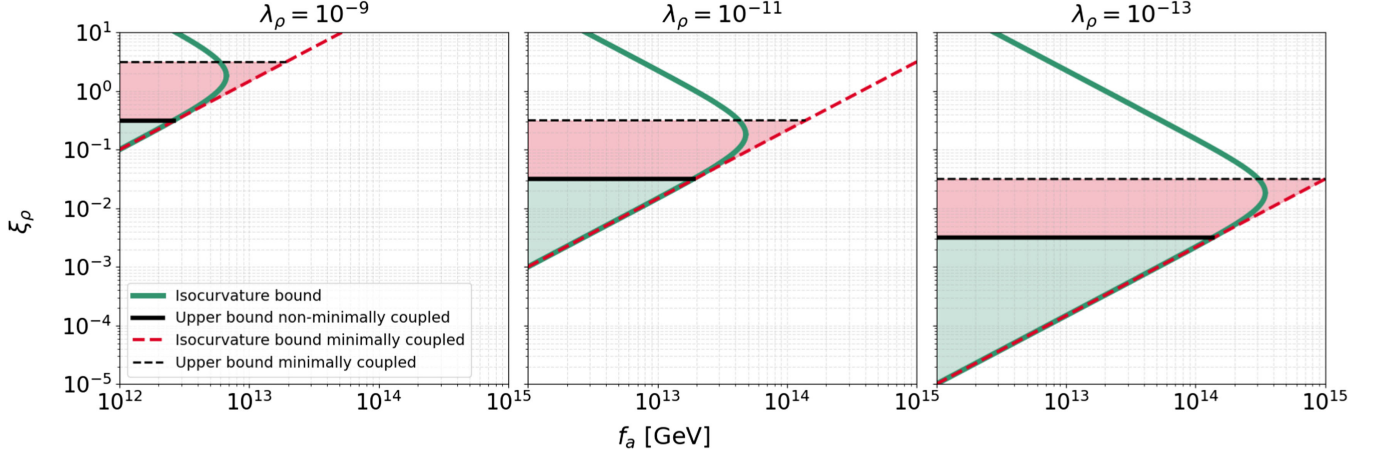


FIG. 2. Constraints on ξ_ρ in Starobinsky inflation as a function of f_a , for $H = 10^{13}$ GeV and some choices of λ_ρ . The green line represents the lower bound on ξ_ρ for which the inflationary decay constant (42) fulfills the isocurvature bound (9). The red curve shows the would-be lower bound on ξ_ρ if the inflaton were minimally coupled, which is found by plugging Eq. (16) into Eq. (9). Furthermore, the solid black line represents the upper bound on ξ_ρ from imposing the nonbackreaction condition of Eq. (46) (with $\Omega_0^2 = N = 50$) and the dashed black line corresponds to the would-be upper bound for the minimally coupled case. In the green region, isocurvature bounds are obeyed and the red region would only be viable if the influence of the inflaton nonminimal coupling on axions were neglected. The code used to produce this figure is available at Ref. [67].

$$U_\rho \approx \frac{9\lambda_\rho}{64N^2} \left(\left(1 + \frac{12\xi_\rho^2 H^2}{\lambda_\rho M_P^2} \right) \rho^4 - 2 \left(f_a^2 + \frac{16N\xi_\rho H^2}{\lambda_\rho} \right) \rho^2 + f_a^4 \right), \quad (40)$$

where we used that $H = \sqrt{V/3}/M_P \approx M/2$ and $\chi \sim M_P \sqrt{3/2} \ln(4N/3)$ in Starobinsky inflation (see [29]). Thus, the minimum of the potential is at

$$\rho_{\min} \approx f_a \sqrt{\frac{1 + \frac{16N\xi_\rho H^2}{\lambda_\rho f_a^2}}{1 + \frac{12\xi_\rho^2 H^2}{\lambda_\rho M_P^2}}}, \quad (41)$$

and we arrive at the inflationary decay constant

$$f_a^{(\text{inf})} = \frac{\rho_{\min}}{\Omega} \approx f_a \sqrt{\frac{\frac{3}{4N} + \frac{12\xi_\rho H^2}{\lambda_\rho f_a^2}}{1 + \frac{12\xi_\rho^2 H^2}{\lambda_\rho M_P^2}}}. \quad (42)$$

We can approximate

$$f_a^{(\text{inf})} \simeq \begin{cases} \sqrt{\frac{3}{4N}} f_a & \text{for } \xi_\rho \ll \frac{\lambda_\rho f_a^2}{NH^2} \\ \sqrt{\frac{12\xi_\rho H^2}{\lambda_\rho}} & \text{for } \frac{\lambda_\rho f_a^2}{NH^2} \ll \xi_\rho \ll \sqrt{\lambda_\rho} \frac{M_P}{H}, \\ \frac{M_P}{\sqrt{\xi_\rho}} & \text{for } \xi_\rho \gg \sqrt{\lambda_\rho} \frac{M_P}{H} \end{cases} \quad (43)$$

where we assumed $\sqrt{\lambda_\rho} M_P/H \gg \lambda_\rho f_a^2/(NH^2)$.¹⁰

Therefore, if ξ_ρ is too small, we get $f_a^{(\text{inf})} < f_a$ and isocurvature bounds are strengthened, as derived in [29].

In order to identify viable parts in parameter space, we first need to make sure to fulfill the condition of not altering inflation. From Eq. (39), we can read off the leading correction to the first derivative of the potential with respect to χ [first term in last line of Eq. (39)]

$$\frac{\sqrt{\epsilon} - \sqrt{\epsilon}|_{\xi_\rho=0}}{\sqrt{\epsilon}|_{\xi_\rho=0}} \sim \xi_\rho \frac{\rho_{\min}^2}{M_P^2}, \quad (45)$$

where we assumed $f_a^4 \lambda_\rho \ll M^2 M_P^2 N$. This leads to the condition

$$\xi_\rho \frac{\rho_{\min}^2}{M_P^2} \ll 1 \Leftrightarrow \xi_\rho \ll \frac{M_P^2}{N f_a^{(\text{inf})2}}, \quad (46)$$

which is stronger than the estimate (34) since $\sigma^2 \sim NM^2$. We conclude that the third line of Eq. (43) cannot obey Eq. (46), and in the second line, the admissible interval of ξ_ρ shrinks.

¹⁰For large f_a close to M_P , it can be possible to achieve a hierarchy $\sqrt{\lambda_\rho} M_P/H \ll \xi_\rho \ll \lambda_\rho f_a^2/(NH^2)$. Then one gets

$$f_a^{(\text{inf})} \simeq \frac{f_a M_P}{4\xi_\rho H} \sqrt{\frac{\lambda_\rho}{N}}, \quad (44)$$

and evidently $f_a^{(\text{inf})} < f_a$.

In summary, we can alleviate isocurvature bounds if

$$\begin{aligned} \frac{\lambda_\rho f_a^2}{H^2} &\ll \xi_\rho \ll \sqrt{\frac{\lambda_\rho}{N}} \frac{M_P}{H} \\ \Rightarrow f_a^{(\text{inf})} &\approx \sqrt{\frac{12\xi_\rho}{\lambda_\rho}} H \gtrsim f_a. \end{aligned} \quad (47)$$

This result for $f_a^{(\text{inf})}$ coincides with its counterpart in Eq. (27) derived in the Palatini formulation of GR with an identical lower bound on ξ_ρ . Furthermore, the upper bound on ξ_ρ is the same as Eq. (29), derived in Palatini Higgs inflation (taking into account that $\Omega^2 \approx N$). Thus, the additional kinetic mixing, which arises in metric GR [see Eq. (7)], does not play a role.

Plugging the largest admissible ξ_ρ of Eq. (47) into the isocurvature bound (10), we get

$$\lambda_\rho < 10^{-7} \left(\frac{10^{12} \text{ GeV}}{f_a} \right)^{7/3}. \quad (48)$$

This result is very similar to its counterpart (30) in Palatini Higgs inflation. Thus, a very small λ_ρ is required to alleviate isocurvature bounds. Because of the condition $T_{\text{max}} < f_a$ of not restoring PQ symmetry during reheating, the precise value of the largest admissible λ_ρ depends the transition from inflation to radiation dominated expansion.¹¹ In Fig. 2, we show the influence of the inflationary decay constant (42) on isocurvature bounds for different values of λ_ρ .

V. FURTHER CONSTRAINTS

In this paper, we have focused on the avoidance of isocurvature constraints. Of course, this is only a necessary but not a sufficient condition for the phenomenological viability of a given scenario. We refer the reader to [51] for a discussion of further constraints, among which we shall briefly discuss the following two.

A nonminimal coupling to gravity with a parameter $\xi_\rho > 1$ lowers the cutoff scale Λ , beyond which perturbation theory breaks down, below the Planck scale. In metric GR, $\Lambda \sim M_P/\xi_\rho$ [73,74], and so the requirement $T_{\text{max}} < \Lambda$ can lead to an upper bound on ξ_ρ that is stronger than the

conditions discussed thus far.¹² In contrast, Palatini gravity leads to the significantly higher $\Lambda \sim M_P/\sqrt{\xi_\rho}$ [75] (see also [76]). Therefore, it follows from $H \ll \sqrt{\lambda_\rho} M_P/\xi_\rho$ [cf., Eqs. (17), (27)] that $T_{\text{max}} < \sqrt{M_P H} < \lambda_\rho^{1/4} M_P/\sqrt{\xi_\rho} < \Lambda$, i.e., the cutoff scale does not further constrain the viable parameter space in Palatini GR.

During reheating, the axionic decay constant relaxes from $f_a^{(\text{inf})}$ to its late-time value f_a . This can lead to a nonthermal restoration of the PQ-symmetry [77–80] due to nonperturbative effects, and possibly also the formation of problematic topological defects [54,81–86]. Without considering the effect of noncanonical kinetic terms, it was suggested that avoidance of PQ-restoration could lead to a strong bound $\xi_\rho \lesssim 10^2 \lambda_\rho f_a^2/H^2$, although the precise numerical value depends on the model [51]. We expect, however, that noncanonical kinetic terms strongly influence the evolution of ρ and a during reheating, and so a detailed investigation of this phenomenon—and reheating in general—remains to be performed.

VI. CONCLUSION

For scalar fields, the presence of a nonminimal coupling to gravity is arguably more natural than its absence, and such couplings play a central role in many successful inflationary models. However, a nonminimal coupling ξ_σ of the inflaton to gravity inevitably decreases the inflationary value $f_a^{(\text{inf})}$ of the axion decay constant and thus worsens the compatibility with isocurvature bounds [29]. Conversely, for a minimally coupled inflaton, a nonminimal coupling ξ_ρ of the radial PQ field can increase $f_a^{(\text{inf})}$ and thereby relax isocurvature constraints [51]. In this paper, we have combined these two effects and identified the conditions under which each of them dominates.

We have shown that ξ_σ reduces the maximal viable value of ξ_ρ , primarily due to the requirement of not significantly modifying the derivative of the inflationary potential. This in turn suppresses the maximal enhancement of the inflationary decay constant $f_a^{(\text{inf})}/f_a$ and therefore reduces the parameter region in which isocurvature bounds can be alleviated. Remarkably, whenever ξ_ρ succeeds in alleviating isocurvature constraints, the same expression (2) for $f_a^{(\text{inf})}$ as in the case of a minimally coupled inflaton still holds (cf., Table I). Furthermore, these findings are largely insensitive to the formulation of GR, although some quantitative features in metric GR remain model dependent.

We have explored these effects in two concrete inflationary scenarios—Palatini Higgs inflation and Starobinsky

¹¹In Starobinsky inflation, reheating is not instantaneous but proceeds more slowly, where details depend on the precise coupling to matter fields (see, e.g., [51,68–72]). As a result, the maximal temperature T_{max} does not saturate the bound (19). On the other hand, $H \sim 10^{13}$ GeV is larger in the Starobinsky scenario as compared to Palatini Higgs inflation, and so the right-hand side of (19) evaluates to a larger number. Therefore, we expect the resulting bounds on λ_ρ to roughly be on the same order in both models.

¹²In [51], the more restrictive condition $\sqrt{M_P H} \lesssim \Lambda$ was imposed. This coincides with our requirement only if reheating can be approximated as instantaneous, i.e., the bound (19) on T_{max} is saturated.

TABLE I. Summary of bounds on ξ_ρ for different models. The lower bound comes from requiring $f_a^{(\text{inf})} > f_a$. The upper bound is necessary to ensure that backreaction on the inflaton is avoided.

Model	$\xi_{\rho,\text{min}}$	$\xi_{\rho,\text{max}}$	$f_a^{(\text{inf})}$
$\xi_\sigma = 0$	$\frac{\lambda_\rho f_a^2}{H^2}$	$\sqrt{\lambda_\rho} \frac{M_P}{H}$	$\sqrt{\frac{12\xi_\rho H^2}{\lambda_\rho}}$
Palatini		$\sqrt{\lambda_\rho} \frac{M_P}{H}$ ^a	
Palatini Higgs		$\frac{\sqrt{\lambda_\rho} M_P}{\Omega_0 H}$	
Metric Starobinsky			

^aNote that the upper bound for Palatini inflation is necessary but not sufficient, and once we specify the potential we may get a stricter upper bound, as shown for Palatini Higgs inflation.

inflation—with all our main findings summarized in Table I. In both cases, a nonminimal coupling $\xi_\rho \sim 10^{-1}$ can alleviate isocurvature bounds, albeit at the cost of requiring a small self-coupling $\lambda_\rho \lesssim 10^{-9}$. It would be very interesting to perform a comprehensive phenomenological parameter scan of these models, in analogy to the analysis of [51].

Furthermore, understanding the role of nonminimal couplings (and the resulting noncanonical kinetic terms) during reheating is an important next step, as it could significantly impact the relevance of nonperturbative effects.

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DATA AVAILABILITY

No data were created or analyzed in this study.

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